Two Model-Independent Results for the Momentum Dependence of ρ-ω Mixing

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Abstract

Two model-independent results on the momentum-dependence of ρ-ω mixing are described. First, an explicit choice of interpolating fields for the vector mesons is displayed for which both the mixing in the propagator and the isospin-breaking at the nucleon-vector meson vertices (and hence also the one-vector-meson-exchange contribution to NN charge symmetry breaking) vanish identically at $q^2 = 0$. Second, it is shown, using the constraints of unitarity and analyticity on the spectral function of the vector meson propagator, that there is no possible choice of interpolating fields for the ρ^0, ω^0 mesons such that, with the $\rho\omega$ element of the propagator defined by $\Delta^{\rho\omega}_{\mu\nu}(q^2) = (g_{\mu\nu} - q_\mu q_\nu/q^2)\theta(q^2)/(q^2 - m_\rho^2)(q^2 - m_\omega^2)$, $\theta(q^2)$ is independent of momentum. It follows that the standard treatment of charge symmetry breaking in few-body systems cannot be interpreted as arising from any realizable effective meson-baryon Lagrangian and must, therefore, be considered purely phenomenological in content.

In standard meson-exchange models of few-body systems, isospin-breaking meson-meson mixing plays an important role in generating contributions to few-body charge-symmetry-violating (CSV) observables. Among these contributions, those associated with ρ-ω mixing have, traditionally, been thought to be rather well-determined, the mixing matrix element being taken (under the somewhat questionable assumption of the absence of direct $\omega^0 \to \pi\pi$ contributions) to be directly measured in the region of the ρ-ω interference shoulder in $e^+e^- \to \pi^+\pi^-$. Using this "extracted" value (i.e. assuming no "direct" $\omega^0 \to \pi\pi$ coupling, where ω^0 is the pure isospin zero ω state) one obtains significant contributions to a number of observables, in particular, the bulk of the non-Coulombic A=3 binding energy difference, significant contributions to the np asymmetry at 183 MeV and non-negligible contributions to the difference of nn and pp scattering lengths and the np asymmetry at 477 MeV [1–9]. This phenomenological success has, however, been recently called into question by the suggestion that such mixing matrix elements must, in general, be expected to be rather momentum-dependent [10–22]. If this were, indeed, the case then, even assuming the interference in $e^+e^- \to \pi^+\pi^-$ in the vicinity of $q^2 \sim m_\omega^2$ were to be dominated by the ρ-ω mixing contribution, the experimental input would not determine the value of the mixing for $q^2 < 0$, where it is needed in few-body CSV calculations.

A significant problem, which has considerably complicated the discussion of this issue in the literature, is the dependence of off-shell Green functions (such as the off-shell propagator) on the choice of interpolating fields. As is well-known, there is no unique choice of fields to represent, eg., the ρ, ω mesons: given a particular choice $\{\rho, \omega\}$ and a corresponding effective Lagrangian, $L_{\rm eff}[\rho, \omega, \cdots]$ (where $\cdots$ represents all other fields), one may define $\rho = \rho' F(\rho')$ and $\omega = \omega' G(\omega')$, with $F(0) = G(0) = 1$ and $L'_{\rm eff}[\rho', \omega', \cdots] \equiv L_{\rm eff}[\rho' F(\rho'), \omega' G(\omega'), \cdots]$. For any such field redefinition, the $\{\rho, \omega, L_{\rm eff}[\rho, \omega, \cdots]\}$ and $\{\rho', \omega', L_{\rm eff}[\rho', \omega', \cdots]\}$ theories have exactly the same S-matrix elements [23,24] and hence are physically equivalent. The Green functions of the two theories, however, are not, in general, the same (for useful pedagogical examples of this statement see, eg., Refs. [25,26]). Thus, the off-shell dependence of the ρ^0-ω^0 element of the vector meson propagator matrix,

$$
\begin{aligned}
\Delta^{\rho\omega}_{\mu\nu}(q^2) &\equiv i\int d^4q\,\exp(iq.x)\ <0|T(\rho^0_\mu(x)\omega^0_\nu(0))|0> \\
&\equiv \left(g_{\mu\nu}-\frac{q_\mu q_\nu}{q^2}\right)\Delta^{\rho\omega}(q^2) \\
&\equiv \left(g_{\mu\nu}-\frac{q_\mu q_\nu}{q^2}\right)\frac{\theta(q^2)}{(q^2-m_\rho^2)(q^2-m_\omega^2)} \qquad (1)
\end{aligned}
$$

will, in general, be changed when one makes a new choice of ρ^0, ω^0 interpolating fields. One may readily display interpolating field choices for which $\theta(q^2)$ necessarily vanishes at $q^2=0$ [18] (and hence is obviously q^2-dependent), eg.,

$$
\begin{aligned}
\rho^0_\mu &= \frac{g_\rho}{\hat{m}_\rho^2}V^\rho_\mu \\
\omega^0_\mu &= \frac{g_\omega}{\hat{m}_\omega^2}V^\omega_\mu \qquad (2)
\end{aligned}
$$

where $\hat{m}_{\rho,\omega}$ are the ρ, ω masses, $V^\rho_\mu=(\bar{u}\gamma_\mu u-\bar{d}\gamma_\mu d)/2$, $V^\omega_\mu=(\bar{u}\gamma_\mu u+\bar{d}\gamma_\mu d)/6$, and g_ρ, g_ω are the usual vector meson decay constants, defined by $<0|V^{\rho,\omega}_\mu|\rho,\omega(q,\epsilon^\lambda)>\equiv \hat{m}^2_{\rho,\omega}\epsilon^\lambda_\mu/g_{\rho,\omega}$, with ϵ^λ the polarization vector. However, given the freedom of field redefinition, this is not enough to exclude the possibility that, for some other field choice, $\theta(q^2)$ might turn out to be q^2-independent. The field redefinitions necessary to produce this effect would then shift the q^2-dependence from the propagator into the vertices in the new effective Lagrangian. As pointed out by Cohen and Miller [27], this raises the possibility that the standard treatment described above might simply correspond to a different interpolating field choice than those of the other treatments, one for which $\theta(q^2)$ is q^2-independent and the CSV vertices, simultaneously, happen to approximately vanish. There is, at present, nothing to rule out this scenario. Given the wide range of field choices made possible by the freedom of field redefinition, it would seem unlikely that one could make further progress. However, we will see below that one can in fact show that (1) there exist interpolating field choices for the vector mesons for which the full, single vector meson exchange contribution to NN CSV vanishes at $q^2=0$ and (2) certain general constraints, associated with unitarity and analyticity, which must be satisfied for all choices of interpolating field, exclude the possibility of finding interpolating fields for which $\theta(q^2)$ is constant.

In order to set the context for the first point above, it is useful to begin with a slightly generalized form of an observation first made by Cohen and Miller. This states that there is no *algebraic* distinction between CSV NN interactions (at the one boson exchange level) associated with CSV-vertex and CSV-propagator contributions. The argument required to arrive at this observation runs as follows. To first order in isospin breaking, the CSV contributions to NN scattering associated with vector meson exchange are of two types: (1) those involving one charge symmetry conserving (CSC) and one CSV $\rho^0 NN$ (or $\omega^0 NN$) vertex, combined with a (CSC) ρ^0 (or ω^0) propagator, and (2) those involving CSC $\rho^0 NN$ and $\omega^0 NN$ vertices together with the CSV off-diagonal $\rho^0\omega^0$ element of the vector meson propagator. If we consider the latter contribution, it is (suppressing the Lorentz indices, γ matrices and nucleon spinors, which are inessential to the argument), of the form

$$v_\tau^{(1)}(q^2)\frac{\theta(q^2)}{(q^2-m_\rho^2)(q^2-m_\omega^2)}v^{(2)}(q^2) \tag{3}$$

where $v_\tau^{(1)}$ is the CSC (isovector) $\rho^0 NN$ and $v^{(2)}$ the CSC (isoscalar) $\omega^0 NN$ vertex, and the q^2-dependence of the vertices results from phenomenological form factors, which are supposed to provide a representation of higher order effects in the effective meson-baryon theory. Now imagine that we write $\theta(q^2) = c + [\theta(q^2) - c]$ and use the standard partial fraction decomposition

$$\frac{1}{(q^2-m_\rho^2)(q^2-m_\omega^2)} = \frac{1}{(m_\omega^2-m_\rho^2)}\left[\frac{1}{q^2-m_\omega^2} - \frac{1}{q^2-m_\rho^2}\right] . \tag{4}$$

The expression (3) can then be re-written as

$$\begin{aligned} v_\tau^{(1)}(q^2)\frac{c}{(q^2-m_\rho^2)(q^2-m_\omega^2)}v^{(2)}(q^2) + v_\tau^{(1)}(q^2)\frac{b(q^2)}{(q^2-m_\omega^2)}v^{(2)}(q^2) \\ -v_\tau^{(1)}(q^2)\frac{b(q^2)}{(q^2-m_\rho^2)}v^{(2)}(q^2) \end{aligned} \tag{5}$$

where

$$b(q^2) \equiv \frac{\theta(q^2)-c}{m_\omega^2-m_\rho^2} . \tag{6}$$

The interesting observation made by Cohen and Miller is that the expression (5) is of precisely the same form as would result from a combination of three contributions: (1) a CSV

mixed $\rho^0\omega^0$ exchange having constant $\theta(q^2) = c$ and CSC vertices, (2) a CSC ω^0 exchange with CSV vertex $v_\tau^{(1)}(q^2)b(q^2)$ and CSC vertex $v^{(2)}(q^2)$, and (3) a CSC ρ^0 exchange with CSV vertex $v^{(2)}(q^2)b(q^2)$ and CSC vertex $v_\tau^{(1)}(q^2)$. So far this is no more than algebraic manipulation. It can, however, be used to give physical meaning to the standard treatment of few-body CSV if, first, noting that $\theta(m_\omega^2)$ is "measured" in $e^+e^- \to \pi^+\pi^-$, one considers, in the language of the discussion above, $c = \theta(m_\omega^2)$, and, second, having made this choice, is able to argue that the function $v^{(2)}(q^2)b(q^2)$ can be interpreted as a *physical* $\rho^0 NN$ CSV vertex (for this particular choice of c, the second term in Eqn. (5) becomes non-singular at $q^2 = m_\omega^2$ and hence is of the form of a background contribution). Eqn. (5), with $c = \theta(m_\omega^2)$, would then correspond to an effective theory with constant $\theta(q^2)$, in which the additional "vertex-like" terms in (5) are viewed as being only a part of the full contributions associated with the CSV $\rho^0 NN$, $\omega^0 NN$ vertices of the theory. If, combined with the other CSV vertex terms, the net CSV vertex contributions were to be small, one would then have recovered the standard treatment. Given the phenomenological successes of this approach, it is then suggested that alternate approaches which find large q^2-dependence of $\theta(q^2)$ might, if they evaluated the CSV vertices using the same choice of fields, find that the vertex contributions essentially cancelled the effects of the q^2-dependence of $\theta(q^2)$, restoring the standard treatment. We now show, however, that this cannot, in general, be true by giving an explicit choice of ρ^0, ω^0 interpolating fields for which both $\theta(q^2)$ and the CSV vertices vanish at $q^2 = 0$.

Let us consider the simplest (and most natural) field choices for the ρ^0 and ω^0, given by Eqns. (2). The off-diagonal element of the vector meson propagator matrix then becomes, up to a constant, just the current correlator $< 0|T(V_\mu^\rho V_\nu^\omega)|0 >$, for which $\theta(q^2) = 0$ [18]. V_μ^ρ, moreover, is just the third component of the isospin current, and V_ν^ω a linear combination of the hypercharge and baryon number currents, and all three of these currents remain exactly conserved, even in the presence of isospin breaking. As such, the nucleon matrix elements of these currents are uniquely determined at $q^2 = 0$ by the I_3, Y and B values of the

nucleon. In particular, there are no contributions to the $q^2 = 0$ values of $< N'|V^\rho_\mu|N >$ and $< N'|V^\omega_\nu|N >$ at any order in $(m_d - m_u)$ or $\alpha_{\rm EM}$, and hence no CSV vertex contributions at $q^2 = 0$. The full NN CSV contribution due to single ρ^0, ω^0 vector meson exchange graphs of the theory having these interpolating fields thus vanishes at $q^2 = 0$, in contrast to the non-zero contribution obtained in the standard treatment. All CSV at $q^2 = 0$ in this case would then have to be associated with multiple meson and/or heavier meson exchanges and the full single-vector-meson-exchange CSV contribution would actually change sign in going from the timelike to the spacelike region.

It should be stressed that the argument above does not imply that all NN CSV vanishes at $q^2 = 0$, only that associated with single ρ^0, ω^0 exchange. One must, moreover, bear in mind that it is only the full S-matrix, and not the one-boson-exchange contribution thereto, which is independent of interpolating field choice. As such, one might still entertain the weaker hypothesis that, in spite of the above behavior for the interpolating field choice of Eqns. (2), there exists some other field choice for which the standard scenario is realized. This brings us to our second point, which is to show that even this weaker hypothesis is untenable.

To begin, let us be more specific about what would be required to give the algebraic manipulation above a physical, as opposed to simply a phenomenological, meaning, namely that there should be *some* choice of ρ^0, ω^0 interpolating fields for which the contributions having the algebraic form of either the CSV propagator mixing or CSV vertex terms in (5) are actually generated by CSV in the vector meson propagator matrix or vector meson-nucleon vertices of the corresponding effective theory $L_{\rm eff}[\rho^0, \omega^0, \cdots]$. Although, as first noted by Cohen and Miller, the manipulation above demonstrates that there is no *algebraic* distinction between such sources of CSV, the distinction becomes crucial once one wishes to use information extracted from $e^+e^- \rightarrow \pi^+\pi^-$ in the few-body context, since only the CSV-propagator, and not the CSV-vertex, contributions are present, and hence (potentially) extractable from the $e^+e^- \rightarrow \pi^+\pi^-$ experiment. Although, given the freedom of field redefinition, a huge class of possible $L_{\rm eff}$'s exists, and one might despair of obtaining any general

information, valid for all of them, it is presumably the case that, regardless of interpolating field choice, certain general properties, such as unitarity, analyticity and the existence of a spectral representation of the propagator, must be common to all of them. As we will now see, this constraint is enough to rule out the possibility that there exists any choice of interpolating fields for which $\theta(q^2)$ is constant. This is our second result.

Let us consider the off-diagonal element of the scalar propagator function, $\Delta^{\rho\omega}(q^2)$, defined in Eqn. (1). The general form of the corresponding spectral function is known: it has poles at $q^2 = m_\rho^2,\, m_\omega^2$, where $m_\rho^2 = \hat{m}_\rho^2 - i\Gamma_\rho \hat{m}_\rho$ and $m_\omega^2 = \hat{m}_\omega^2 - i\Gamma_\omega \hat{m}_\omega$, and cuts along the positive, real q^2 axis, the first of these beginning at $q^2 = 4m_\pi^2$. $\Delta^{\rho\omega}(q^2)$ is then real for q^2 real and $< 4m_\pi^2$. In what follows we will ignore the width of the ω.

Recall that, from the definition of $\theta(q^2)$ in Eqn. (1), one has

$$\theta(q^2) = (q^2 - m_\rho^2)(q^2 - m_\omega^2)\Delta^{\rho\omega}(q^2) \tag{9}$$

where the vector meson squared-masses are the *complex* pole positions. It is crucial to use this form, rather than that in which the complex pole positions have been replaced by their real parts, in order to make contact with the data from $e^+e^- \to \pi^+\pi^-$, since the extracted value of θ is obtained using a fitting form in which the complex pole locations are explicitly present in the resonant denominators (see Ref. [28]). One can now easily see that $\theta(q^2)$ cannot possibly be constant, at least below $q^2 = 4m_\pi^2$. Indeed, taking the ratio of Eqn. (7) at two different below-threshold values, one has

$$\frac{\theta(q_1^2)}{\theta(q_2^2)} = \left[\frac{q_1^2 - m_\rho^2}{q_2^2 - m_\rho^2}\right]\left[\frac{(q_1^2 - m_\omega^2)\Delta^{\rho\omega}(q_1^2)}{(q_2^2 - m_\omega^2)\Delta^{\rho\omega}(q_2^2)}\right] \ . \tag{8}$$

The second factor on the RHS of Eqn. (8) is (neglecting the ω width) real, the first term complex when $q_1^2 \neq q_2^2$. Thus $\theta(q_1^2) \neq \theta(q_2^2)$ for $q_1^2 \neq q_2^2$ and both below threshold: constancy of $\theta(q^2)$ below threshold is incompatible with the required spectral behavior of the propagator.

One can, in fact, use Eqn. (8), together with the reality of $\Delta^{\rho\omega}(q^2)$ below threshold, to put an "analyticity constraint" on the q^2-variation of $\theta(q^2)$, i.e. to give a lower bound on

the q^2-variation of $\theta(q^2)$ below threshold compatible with the constraints of unitarity and analyticity on the spectral function. Let us re-write Eqn. (8) as

$$\frac{\theta(q_1^2)}{\theta(q_2^2)} - 1 = r(q_1^2, q_2^2)\left[\frac{q_1^2 - m_\rho^2}{q_2^2 - m_\rho^2}\right] - 1 \tag{9}$$

where, as we have seen, $r(q_1^2, q_2^2)$ is real for q_1^2,$q_2^2 < 4m_\pi^2$, if we ignore the ω width. One may then ask, if one is allowed to adjust the form of $\Delta^{\rho\omega}(q^2)$ (and hence $r(q_1^2, q_2^2)$) so as to minimize the magnitude of the RHS of Eqn. (9), how small can the q^2-dependence be made, subject only to the constraint that $r(q_1^2, q_2^2)$ remain real? Note that there is no guarantee that one could actually succeed in finding interpolating fields which realize this lower bound, given that the requirement corresponds to a highly restrictive statement about the form and magnitude of the off-diagonal element of the propagator, and as a result, the actual q^2-variations will, in general, be larger (probably much larger) than those obtained following the procedure just described. What one, however, is guaranteed, is that the actual variation with q^2 for *any* choice of interpolating fields must be greater than that specified by the bound so obtained. Taking the magnitudes of both sides of Eqn. (9), and determining the (real) value of $r(q_1^2, q_2^2)$ which minimizes this magnitude, one finds, for example, a minimum variation in the magnitude of $\theta(q^2)$ of 15% between $q^2 = -1\,\mathrm{GeV}$ and $q^2 = 0$. We stress again that actually reducing the variation to this level is not necessarily possible, and that typical variations (as in the case of the vector current interpolating field choice) will be much greater. Moreover, although the argument above can no longer be implemented for $q^2 > 4m_\pi^2$, the fact that q^2-variation is unavoidable below $q^2 = 4m_\pi^2$ clearly argues for the likelihood of its presence above this point.

The validity of the argument above, of course, rests on the fact that the ρ is not a narrow resonance. If, instead, both the ρ and ω had essentially zero width, then the conclusion could be completely evaded, as is evident from Eqn. (8). The significant effect of including the width of the ρ, has also been stressed recently in Ref. [22], where it is demonstrated that taking the spectral function of $\Delta^{\rho\omega}$ to consist of a sum of constant multiples of the ρ and ω Breit-Wigner resonance forms leads to significant q^2-dependence of $\Delta^{\rho\omega}$. While such a form

for the spectral function is not the most general that would be produced if one considered all possible field redefinitions, (only S-matrix properties, like the pole positions, are independent of the interpolating field choice), it should be noted that the effect on the q^2-dependence of including the ρ width is numerically very large.

Another alternative for evading the argument would be to write, instead of (7),

$$\theta'(q^2) \equiv (q^2 - \hat{m}_\rho^2)(q^2 - \hat{m}_\omega^2)\Delta^{\rho\omega}(q^2) \tag{10}$$

and assume $\theta'(q^2)$ was constant, which is then consistent with $\Delta^{\rho\omega}(q^2)$ being purely real below threshold. Now, however, one can no longer obtain θ' from $e^+e^- \to \pi^+\pi^-$ (which is analyzed using the alternate form having the correct pole locations). $\theta'(q^2)$, moreover, has a pole at $q^2 = m_\rho^2$ and a zero at $q^2 = \hat{m}_\rho^2$, and hence is certainly not constant near $q^2 = \hat{m}_\rho^2$. The resulting rapid variation in the vicinity of $q^2 = \hat{m}_\rho^2$ would also make the connection of the value of θ measured experimentally (even assuming direct $\omega \to \pi\pi$ contributions *can* be neglected) to the values of $\theta'(q^2)$ for q^2 on the real axis below the position of the zero, rather problematic. Finally, one easily sees that, even ignoring the existence of a nearby pole and zero, θ' cannot be constant over the whole of the range required to save the "standard" treatment, since $\Delta^{\rho\omega}$ has a non-zero imaginary part above $q^2 = 4m_\pi^2$, presumably significantly so in the region of the ρ peak, and this means that, if θ' were constant with q^2, $\Delta^{\rho\omega}$ would then also have a non-zero imaginary part, eg. at $q^2 = 0$, incompatible with the requirement that it be real below threshold. This argument can be evaded, again, only when both resonances have zero width (in which case the ρ pole moves up to the real axis, cancelling the factor of $(q^2 - \hat{m}_\rho^2)$, producing a non-zero real part of θ' and an imaginary part which vanishes for all q^2, assuming all spectral strength to be isolated at the ρ and ω poles).

In light of the argument above, it appears that the standard approach to few-body CSV cannot be physically justified. While one might worry about the fate of the phenomenological successes associated with the standard treatment, it is worth noting that, first, this success is based on the potentially dangerous assumption of the neglect of direct $\omega^0 \to \pi\pi$ contributions

to $e^+e^- \rightarrow \pi^+\pi^-$ and, second, a $\rho^0\omega^0$ mixing contribution that is, for example, half the size of that usually employed would, in fact, create no phenomenological problems, except for a somewhat smaller than required non-Coulombic contribution to the $A = 3$ binding energy difference.

Concerning the first point, it should be stressed that, although direct $\omega^0 \rightarrow \pi\pi$ contributions to $e^+e^- \rightarrow \pi^+\pi^-$ are usually neglected, there is actually no good reason to assume that they will be negligible relative to those associated with $\rho^0\omega^0$ mixing. Indeed, both are, in general, non-zero at $\mathcal{O}(m_d - m_u)$, and should, therefore, barring other information, be expected to be comparable in magnitude, as is born out by both a recent QCD sum rule analysis of the vector current correlator $< 0|T(V_\mu^\rho V_\nu^\omega)|0 >$ [21] and a recent calculation of the direct $\omega^0 \rightarrow \pi\pi$ contribution in a model using confining quark propagators and Bethe-Salpeter meson-quark vertices motivated by non-perturbative Schwinger-Dyson equation studies [29]. Without the neglect of direct $\omega^0 \rightarrow \pi\pi$ contributions, however, $e^+e^- \rightarrow \pi^+\pi^-$ cannot be used to give direct information on $\Delta_{\mu\nu}^{\rho\omega}$. Significant direct $\omega^0 \rightarrow \pi\pi$ contributions, as suggested by the studies mentioned above, would then, of course, call the phenomenological successes of the standard approach in question.

Concerning the second point (the non-Coulombic contributions to the $A = 3$ binding energy difference), it should be noted that, as pointed out in Refs. [30–33], it is likely that electromagnetic effects associated with photon-loop, rather than photon-exchange graphs (the former cannot be disentangled from the strong-interaction, QCD effects in a model-independent fashion) play a significant role in CSV in few-body systems, just as, in order to satisfy chiral constraints on the pion electromagnetic self-energies, they must in the pseudoscalar spectrum [30–33].

In summary, it has been shown that (1) there exist interpolating field choices for the vector mesons for which the single-vector-meson-exchange contribution to NN CSV vanishes identically at $q^2 = 0$ and (2) the assumption of a momentum-independent $\theta(q^2)$ for the $\rho^0\omega^0$ propagator is incompatible with the constraints on the spectral function associated with analyticity and unitarity. As such, there can be no choice of ρ^0, ω^0 interpolating fields for

which the standard approach to few-body CSV is physically realizable. Since the standard approach cannot be interpreted as arising from any effective meson-baryon Lagrangian it must, in consequence, be interpreted as being purely phenomenological in nature. Given the extremely strong assumptions required to reduce the q^2-variation to even the level of the analyticity bound discussed above, it seems clear that, if one wishes to neglect the role of these effects in few-body systems, it will be necessary to demonstrate explicitly the existence of interpolating field choices for which the standard assumption is an acceptable approximation.

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